VERIFIABLE TYPE-II SEESAW AND DARK MATTER IN A GAUGED $U(1)_{B-L}$ MODEL

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We propose a gauged $U(1)_{B-L}$ extension of the standard model (SM) to explain simultaneously the light neutrino masses and dark matter (DM). The generation of neutrino masses occurs through a variant of type-II seesaw mechanism in which one of the scalar triplets lies at the TeV scale yet have a large dilepton coupling, which paves a path for probing this model at colliders. The gauging of $U(1)_{B-L}$ symmetry in a type-II seesaw framework introduces $B-L$ anomalies. Therefore we invoke three right handed neutrinos ν_{R_i}(i=1,2,3) with $B-L$ charges -4,-4,+5 to cancel the anomalies. We further show that the lightest one among the three right handed neutrinos can be a viable DM candidate. The stability of DM can be owed to a remnant Z_2 symmetry under which the right handed neutrinos are odd while all other particles are even. We then discuss the constraints on the model parameters from observed DM abundance and the search at direct detection experiments.

I. INTRODUCTION

The Standard Model (SM) of Particle physics works remarkably well in describing the electroweak and strong interactions of fundamental particles of nature. But there are many questions which are unanswered till date. Among them the most popular unsolved problems are the identity of Dark Matter (DM) and the origin of neutrino masses. The cosmological observations like the galaxy rotation curve, gravitational lensing and large scale structure of the Universe provide the evidences towards the existence of DM [1, 2]. But we do not have much information about microscopic properties of DM apart from its relic density, which is precisely measured by the WMAP [3] and PLANCK [4] to be $\Omega_{\rm DM}h^2 = 0.120 \pm 0.001$.

Initially the neutrinos were thought to be massless particles due to lack of experimental evidences. But the neutrino oscilation experiments [5–7] confirmed that they are massive but tiny. Assuming that the neutrinos are Majorana ($\triangle L = 2$), their sub-eV masses are best understood by the dimension five operator: $\mathcal{O}_5 = \frac{LLHH}{\Lambda}$, where L and H are the lepton and Higgs doublets of the SM and Λ is the scale of new physics [8, 9]. After electroweak phase transition, we get sub-eV

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neutrino masses $M_\nu = \mathrm{O}(\frac{\langle H\rangle^2}{\Lambda}) \simeq 0.1$ eV, for $< H >= 10^2$GeV and $\Lambda = 10^{14}$GeV. Seesaw mechanisms: Type-I [10–13], Type-II [14–18] and Type-III [19] are the various UV completed realizations of this dimension five operator. In the type-I seesaw one introduces heavy singlet RHNs while in case of type-II and type-III, one introduces a triplet scalar(Δ) of hypercharge 2 and triplet fermions of hypercharge 0 respectively.

In the conventional type-II seesaw, the relevant terms in the Lagrangian which violates lepton number by two units are $f_{\alpha\beta}\Delta L_\alpha L_\beta + \mu\Delta^\dagger HH$, where Δ does not acquire an explicit vacuum expectation value(vev). However, the electro-weak phase transition induces a small vev of Δ as: $\langle\Delta\rangle = -\frac{\mu\langle H\rangle^2}{M_\Delta^2}$. Thus for $\mu \sim M_\Delta \sim 10^{14}$GeV, one can get $M_\nu = f\langle\Delta\rangle \simeq f\frac{\langle H\rangle^2}{M_\Delta}$ of order O(0.1)eV for $f \sim 1$. The only drawback in this case is that the mass scale of the scalar triplet is much larger than the energy attainable at current generation colliders. Hence such models lack falsifiability.

Alternatively one can introduce two scalar triplets: Δ and ξ with $M_\Delta \sim 10^{14}$GeV and $M_\xi \sim$ TeV $<< M_\Delta$[20] [1]. If one imposes an additional $B-L$ gauge symmetry [21], then the relevant terms in the Lagrangian are: $\mu\Delta^\dagger HH + f\xi LL + y\Phi_{B-L}^2\Delta^\dagger\xi$, where Φ_{B-L} is the scalar field responsible for $B-L$ symmetry breaking. At TeV scales Φ_{B-L} acquires a vev and break $B-L$ symmetry. Moreover $\langle\Phi_{B-L}\rangle$ generates a small mixing between Δ and ξ of the order $\theta \sim \frac{\langle\Phi_{B-L}\rangle^2}{M_\Delta^2} \simeq 10^{-18}$. This implies that ξLL coupling can be large while ξ's coupling with Higgs doublet is highly suppressed. Since Δ mass is super heavy, it gets decoupled from the low energy effective theory. On the other hand, ξ can be at TeV scale with large dilepton coupling. As a result the same sign dilepton signature of ξ can be studied at colliders [23–30].

The gauging of $U(1)_{\rm B-L}$ symmetry in a type-II seesaw framework introduces non-trivial gauge and gravitational anomalies. With the SM particle content all triangle anomalies are zero except for $\sum[U(1)_{B-L}]^3 = 3$ and $\sum[Grav.]^2 \times [U(1)_{B-L}] = 3$. These anomalies can be cancelled by introducing new fermions in such a way that sum of their $B-L$ quantum numbers is -3. In this paper we introduce three right handed neutrinos ν_{R_i} $(i = 1, 2, 3)$ with $U(1)_{B-L}$ charges $-4, -4, +5$ respectively, such that $\sum_{i=1}^3 (Y_{B-L})_i = -3$, to make the theory anomaly free [31–38]. Interestingly, one of the three right handed neutrions can be a viable candidate of DM. The stability of the DM candidate can be guaranteed by a remanant Z_2 symmetry of the original $U(1)_{B-L}$. Under the Z_2 discrete symmetry ν_{R_i} $(i = 1, 2, 3)$ are odd while all other particles are even. Thus without imposing any additional discrete symmetry we can explain the DM as well as smallness of the neutrino masses.

[1] A modified double type-II seesaw with TeV scale scalar triplet is also proposed in ref. [22]

The paper is organised as follows. In section 2, we briefly discuss the gauge and gravitaional anomalies and hence the anomaly free conditions of a gauged $U(1)_{\rm B-L}$ extension of a type-II seesaw model. In section 3, we describe the proposed model, the scalar masses and mixing and the neutrino mass generation at TeV scale through a variant of type-II seesaw. We then discuss how the particles introduced for anomaly cancellation become viable DM candidate and study the relic density and it's compatibility with the direct detection experiments in section 4. We also briefly discuss the collider signatures of the model in section 5 and finally conclude in section 6.

II. ANOMALIES IN A GAUGED $U(1)_{\rm B-L}$ EXTENSION OF A TYPE-II SEESAW

Within the SM, $U(1)_{\rm B-L}$ is happend to be an accidental global symmetry. However, the gauged $U(1)_{\rm B-L}$ extension of the SM is not anomaly free. Among all the anomalies arising from the triangle diagrams involving the gauge currents, except $\mathcal{A}[U(1)^3_{B-L}]$ and $\mathcal{A}[(Gravity)^2 \times U(1)_{B-L}]$, all others are trivial. Here $\mathcal{A}$ stands for the anomaly coefficient which in a chiral gauge theory is given by [39]:

$$\mathcal{A} = Tr[T_a[T_b, T_c]_+]_R - Tr[Ta[T_b, T_c]_+]_L \,, \tag{1}$$

where the T denotes the generators of the gauge groups and R and L represents the interactions of right and left chiral fermions with the gauge bosons.

Gauging of $U(1)_{\rm B-L}$ symmetry within the SM lead to the following triangle anomalies:

$$\mathcal{A}_1[U(1)^3_{B-L}] = 3$$

$$\mathcal{A}_2[(Gravity)^2 \times U(1)_{B-L}] = 3\,. \tag{2}$$

If three right handed neutrinos, each of having $B-L$ charge -1, are added to the SM, then they result in $\mathcal{A}_1[U(1)^3_{B-L}] = -3$ and $\mathcal{A}_2[(Gravity)^2 \times U(1)_{B-L}] = -3$ which lead to cancellation of above mentioned gauge anomalies. This is the most natural choice to make $U(1)_{B-L}$ models anomaly free. However we can have alternative ways of costructing anomaly free versions of $U(1)_{B-L}$ extension of the SM [31–38]. In particular, three right handed neutrinos with exotic $B-L$ charges -4,-4,+5 can also give rise to vanishing $B-L$ anomalies as follows.

$$\mathcal{A}_1[U(1)^3_{B-L}] = \mathcal{A}_1^{SM}[U(1)^3_{B-L}] + \mathcal{A}_1^{New}[U(1)^3_{B-L}] = 3 + [(-4)^3 + (-4)^3 + (5)^3] = 0$$

$$\begin{aligned}\mathcal{A}_2[(Gravity)^2 \times U(1)_{B-L}] &= \mathcal{A}_2^{SM}[(Gravity)^2 \times U(1)_{B-L}] + \mathcal{A}_2^{New}[(Gravity)^2 \times U(1)_{B-L}] \\ &= 3 + [(-4) + (-4) + (5)] = 0\end{aligned} \tag{3}$$

In a type-II seesaw framework, the SM is usually extended with a triplet scalar of hypercharge 2. In this case, gauging of $U(1)_{\rm B-L}$ symmetry does not lead to any new anomalies apart from the mentioned above. Therefore, in what follows, we consider a type-II seesaw framework with gauged $U(1)_{\rm B-L}$ symmetry, where the $B-L$ anomalies are cancelled by the introduction of three right handed neutrinos with exotic $B-L$ charges -4,-4,+5. The model thus proposed explains the origin of neutrino mass and DM in a minimal set-up.

III. THE COMPLETE MODEL

We study a variant of type-II seesaw model based on the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$, where B and L stands for the usual baryon and lepton numbers, respectively. Two triplet (under $SU(2)_L$) scalars Δ and ξ:

$$\Delta = \begin{pmatrix} \frac{\delta^+}{\sqrt{2}} & \delta^{++} \\ \delta^0 & -\frac{\delta^+}{\sqrt{2}} \end{pmatrix} \quad \text{and} \quad \xi = \begin{pmatrix} \frac{\xi^+}{\sqrt{2}} & \xi^{++} \\ \xi^0 & -\frac{\xi^+}{\sqrt{2}} \end{pmatrix} \tag{4}$$

with $M_\Delta \sim 10^{14}$ GeV and $M_\xi \sim$ TeV $<< M_\Delta$ are introduced, where the $B-L$ charges of Δ and ξ are 0 and 2 respectively. As discussed in the previous section, the additional $U(1)_{B-L}$ gauge symmetry introduces $B-L$ anomalies in the theory. To cancel these $B-L$ anomalies we introduce three right handed neutrinos ν_{R_i} $(i = 1, 2, 3)$, where the $B-L$ charges of ν_{R_1}, ν_{R_2}, ν_{R_3} are -4, -4, +5 respectively as shown in Eq.3. Note that such unconventional $B-L$ charge assignment of the ν_{R_i} $(i = 1, 2, 3)$ forbids their Yukawa couplings with the SM particles. We also introduce three singlet scalars: Φ_{B-L}, Φ_{12} and Φ_3 whose $B-L$ charges are given by: -1, +8, -10 respectively. As a result Φ_{12} and Φ_3 couples to $\nu_{R_{1,2}}$ and ν_{R_3} respectively through Yukawa terms. The vev's of Φ_{12} and Φ_3 provides Majorana masses to the right handed neutrinos, while the vev of Φ_{B-L} provides a small mixing between Δ and ξ. We will show later that the mixing between Δ and ξ plays a key role in generating sub-eV masses of neutrinos. The particle content and their charge assignments are listed in Table.I.

The Lagrangian involving the new fields consistent with the extended symmetry can be written as:

$$\begin{aligned}\mathscr{L} \supset\ & \overline{\nu_{R_a}}\, i\gamma^\mu D_\mu \nu_{R_b} + \overline{\nu_{R_3}} i\gamma^\mu D_\mu \nu_{R_3} + \left|D_\mu X\right|^2 \\ & + Y^\xi_{ij}\ \overline{L^c_i}\ i\tau_2\, \xi\, L_j + Y_{ab}\ \Phi_{12}\ \overline{(\nu_{R_a})^c}\ \nu_{R_b} + Y_{33}\ \Phi_3\ \overline{(\nu_{R_3})^c}\ \nu_{R3} \\ & + Y_{a3} \overline{(\nu_{R_a})^c}\ \Phi_{_{B-L}} \nu_{R_3} + \text{h.c.} - \text{V}(H, \Delta, \xi, \Phi_{B-L}, \Phi_{12}, \Phi_3)\end{aligned} \tag{5}$$

Fields	$SU(3)_c$	$SU(2)_L$	$U(1)_Y$	$U(1)_{B-L}$
$\nu_{R_{1,2}}$	1	1	0	-4
ν_{R_3}	1	1	0	5
Δ	1	3	2	0
ξ	1	3	2	2
Φ_{B-L}	1	1	0	-1
Φ_{12}	1	1	0	+8
Φ_3	1	1	0	-10

TABLE I: New particles and their quantum numbers under the imposed gauge symmetry.

where

$$D_\mu = \Big(\partial_\mu + i\ g_{_{B-L}} Y_{B-L} (Z_{B-L})_\mu\Big)\,.$$

The g_{B-L} is the gauge coupling associated with $U(1)_{B-L}$ and Z_{B-L} is the corresponding gauge boson. Here, $X = \Phi_{B-L}, \Phi_{12}, \Phi_3$. The i, j runs over 1, 2, 3 and a, b runs over 1,2.

The scalar potential of the model can be written as:

$$\begin{aligned}
\mathrm{V}(H, \Delta, \xi, \Phi_{B-L}, \Phi_{12}, \Phi_3) = & -\mu_H^2 H^\dagger H + \lambda_H (H^\dagger H)^2 + M_\Delta^2 \Delta^\dagger \Delta + \lambda_\Delta (\Delta^\dagger \Delta)^2 \\
& + M_\xi^2 \xi^\dagger \xi + \lambda_\xi (\xi^\dagger \xi)^2 - \mu^2_{\Phi_{B-L}} \Phi^\dagger_{B-L} \Phi_{B-L} + \lambda_{\Phi_{B-L}} (\Phi^\dagger_{B-L} \Phi_{B-L})^2 \\
& - \mu^2_{\Phi_{12}} \Phi^\dagger_{12} \Phi_{12} + \lambda_{\Phi_{12}} (\Phi^\dagger_{12} \Phi_{12})^2 - \mu^2_{\Phi_3} \Phi_3^\dagger \Phi_3 + \lambda_{\Phi_3} (\Phi_3^\dagger \Phi_3)^2 \\
& + \lambda_{H\Delta} (H^\dagger H)(\Delta^\dagger \Delta) + \lambda_{H\xi} (H^\dagger H)(\xi^\dagger \xi) \\
& + \lambda_{H\Phi_{B-L}} (H^\dagger H)(\Phi^\dagger_{B-L} \Phi_{B-L}) + \lambda_{H\Phi_{12}} (H^\dagger H)({\Phi_{12}}^\dagger \Phi_{12}) \\
& + \lambda_{H\Phi_3} (H^\dagger H)(\Phi_3^\dagger \Phi_3) + \lambda_{\Delta\xi} (\Delta^\dagger \Delta)(\xi^\dagger \xi) + \lambda'_{\Delta\xi} (\Delta^\dagger \xi)(\xi^\dagger \Delta) \\
& + \lambda_{\Delta\Phi_{B-L}} (\Delta^\dagger \Delta)(\Phi^\dagger_{B-L} \Phi_{B-L}) + \lambda_{\Delta\Phi_{12}} (\Delta^\dagger \Delta)({\Phi_{12}}^\dagger \Phi_{12}) \\
& + \lambda_{\Delta\Phi_3} (\Delta^\dagger \Delta)(\Phi_3^\dagger \Phi_3) + \lambda_{\xi\Phi_{B-L}} (\xi^\dagger \xi)(\Phi^\dagger_{B-L} \Phi_{B-L}) \\
& + \lambda_{\xi\Phi_{12}} (\xi^\dagger \xi)({\Phi_{12}}^\dagger \Phi_{12}) + \lambda_{\xi\Phi_3} (\xi^\dagger \xi)(\Phi_3^\dagger \Phi_3) \\
& + \lambda_{\Phi_{B-L}\Phi_{12}} (\Phi^\dagger_{B-L} \Phi_{B-L})({\Phi_{12}}^\dagger \Phi_{12}) + \lambda'_{\Phi_{B-L}\Phi_{12}} (\Phi^\dagger_{B-L} \Phi_{12})({\Phi_{12}}^\dagger \Phi_{B-L}) \\
& + \lambda_{\Phi_{B-L}\Phi_3} (\Phi^\dagger_{B-L} \Phi_{B-L})(\Phi_3^\dagger \Phi_3) + \lambda'_{\Phi_{B-L}\Phi_3} (\Phi^\dagger_{B-L} \Phi_3)(\Phi_3^\dagger \Phi_{B-L}) \\
& + \lambda_{\Phi_{12}\Phi_3} ({\Phi_{12}}^\dagger \Phi_{12})(\Phi_3^\dagger \Phi_3) + \lambda'_{\Phi_{12}\Phi_3} ({\Phi_{12}}^\dagger \Phi_3)(\Phi_3^\dagger \Phi_{12}) + \mu \Delta^\dagger H H \\
& + y \Phi^2_{B-L} \Delta^\dagger \xi + Y (\Phi^\dagger_{B-L})^2 \Phi_3 \Phi_{12} + h.c.
\end{aligned} \tag{6}$$

It is worth mentioning that the mass squared terms of Δ and ξ are chosen to be positive so that only the neutral components of H,Φ_{12},Φ_{B-L} and Φ_3 acquire non-zero vevs. These vevs are given as follows:

$$\langle H\rangle = \frac{v}{\sqrt{2}}\begin{pmatrix}0\\1\end{pmatrix}, \langle\Phi_{12}\rangle = \frac{v_{12}}{\sqrt{2}}, \langle\Phi_{B-L}\rangle = \frac{v_{B-L}}{\sqrt{2}}, \langle\Phi_3\rangle = \frac{v_3}{\sqrt{2}}.$$

However, after electroweak phase transition, Δ and ξ acquire induced vevs:

$$\langle\Delta\rangle = \frac{u_\Delta}{\sqrt{2}}\begin{pmatrix}0 & 0\\1 & 0\end{pmatrix} \text{ and } \langle\xi\rangle = \frac{u_\xi}{\sqrt{2}}\begin{pmatrix}0 & 0\\1 & 0\end{pmatrix}$$

A. Neutrino Mass

Until Φ_{B-L} gains a vev, there is no mixing between Δ and ξ. Once Φ_{B-L} acquires a vev, say v_{B-L} at TeV scale, it allows a mixing between Δ and ξ through $yv_{B-L}^2\Delta^\dagger\xi$. After electroweak symmetry breaking Δ gets an induced vev, similarly as in the case of traditional type-II seesaw. Since $B-L$ quantum number of Δ is zero, it does not generate Majorana masses of neutrinos. The induced vev acquired by Δ after EW phase transition is given by

$$\langle\Delta\rangle = u_\Delta = -\frac{\mu v^2}{\sqrt{2}(2M_\Delta^2 + \lambda_{H\Delta}v^2 + \lambda_{\Phi_{B-L}\Delta}v_{B-L}^2 + \lambda_{\Phi_{12}\Delta}v_{12}^2 + \lambda_{\Phi_3\Delta}v_3^2)}. \tag{7}$$

The vev of Δ is required to satisfy:

$$\rho \equiv \frac{M_W^2}{M_z^2\cos^2\theta} = \frac{1+2x^2}{1+4x^2} \approx 1 \tag{8}$$

where $x = u_\Delta/v$. The above constraint implies that $|u_\Delta| < \mathcal{O}(1)$ GeV. Since ξ mixes with Δ after $U(1)_{B-L}$ breaking, it also acquire an induced vev after EW symmetry breaking and is given by

$$\langle\xi\rangle = u_\xi = -\frac{yv_{B-L}^2}{2(2M_\xi^2 + \lambda_{\xi\Phi_{B-L}}v_{B-L}^2 + \lambda_{\xi H}v^2 + \lambda_{\xi\Phi_{12}}v_{12}^2 + \lambda_{\xi\Phi_3}v_3^2)}u_\Delta. \tag{9}$$

If we assume $yv_{B-L}^2 = M_\xi^2 \sim \lambda_{\xi\Phi_{B-L}}v_{B-L}^2 \sim \lambda_{\xi H}v^2 \sim \lambda_{\xi\Phi_{12}}v_{12}^2 \sim \lambda_{\xi\Phi_3}v_3^2$, then we get $u_\Delta \simeq u_\xi$, even if there is orders of magnitude difference in the masses of ξ and Δ. Here we additionally assume that $M_\Delta >> v_{\rm B-L}, v_{12} > M_\xi, v_3, v$. After integrating out the heavy degrees of freedom in the Feynman diagram given in Fig.1, we get the Majorana mass matrix of the light neutrinos to be

$$(M_\nu)_{ij} = Y_{ij}^\xi u_\xi = Y_{ij}^\xi \frac{yv_{B-L}^2}{2(2M_\xi^2 + \lambda_{\xi\Phi_{B-L}}v_{B-L}^2 + \lambda_{\xi H}v^2 + \lambda_{\xi\Phi}v_{12}^2 + \lambda_{\xi\Phi_3}v_3^2)}u_\Delta. \tag{10}$$

As $u_\Delta \simeq u_\xi \lesssim \mathcal{O}(1)$ GeV, we get sub-eV neutrino masses as required by the oscillation experiments. Here it is worth mentioning that the mixing between the super heavy triplet scalar Δ and the TeV scale scalar triplet ξ gives rise to the neutrino mass and it is clear from Fig. 1 that ξ strongly couples to leptons while it's coupling with SM Higgs doublet is highly suppressed because of large M_Δ.

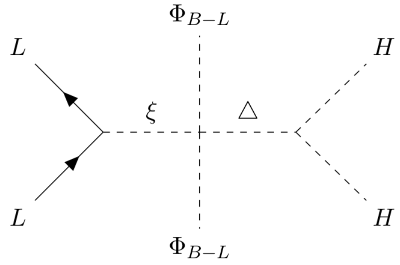

FIG. 1: Generation of neutrino mass through the modified Type-II seesaw.

B. Scalar Masses and Mixing

We parametrize the neutral scalars as:

$$H^0 = \frac{v + h + i\,S}{\sqrt{2}} \quad , \quad \Phi_{B-L} = \frac{v_{B-L} + \phi_{B-L} + iP}{\sqrt{2}}$$

$$\Phi_{12} = \frac{v_{12} + \phi_{12} + i\,Q}{\sqrt{2}} \quad , \quad \Phi_3 = \frac{v_3 + \phi_3 + i\,R}{\sqrt{2}}$$

$$\delta^0 = \frac{u_\Delta + \delta + i\eta}{\sqrt{2}} \quad , \quad \xi^0 = \frac{u_\xi + \xi + i\rho}{\sqrt{2}} \tag{11}$$

The physical mass square terms of the scalars are then given by

$$\begin{aligned}
M_h^2 &\simeq 2\lambda_H v^2 + \frac{\lambda_{H\Phi_{12}} v_{12}^2}{2} + \frac{\lambda_{H\Phi_3} v_3^2}{2} + \frac{\lambda_{H\Phi_{B-L}} v_{B-L}^2}{2} + \frac{\mu u_\Delta}{2} \\
M_\delta^2 &\simeq -\frac{\mu v^2}{4\sqrt{2} u_\Delta} - \frac{y v_{B-L}^2 u_\xi}{8 u_\Delta} \\
M_\xi^2 &\simeq -\frac{y v_{B-L}^2 u_\Delta}{8 u_\xi} \\
M_{\phi_{12}}^2 &\simeq 2\lambda_{\Phi_{12}} v_{12}^2 + \frac{\lambda_{H\Phi_{12}} v^2}{2} + \frac{\lambda_{\Phi_{12}\Phi_3} v_3^2}{2} \\
&+ \frac{\lambda'_{\Phi_{12}\Phi_3} v_3^2}{2} + \frac{\lambda_{\Phi_{12}\Phi_{B-L}} v_{B-L}^2}{2} + \frac{\lambda'_{\Phi_{12}\Phi_{B-L}} v_{B-L}^2}{2} + \frac{Y v_{B-L}^2 v_3}{8 v_{12}} \\
M_{\phi_3}^2 &\simeq 2\lambda_{\Phi_3} v_3^2 + \frac{\lambda_{H\Phi_3} v^2}{2} + \frac{\lambda_{\Phi_{12}\Phi_3} v_{12}^2}{2} \\
&+ \frac{\lambda'_{\Phi\Phi_3} v_{12}^2}{2} + \frac{\lambda_{\Phi_3\Phi_{B-L}} v_{B-L}^2}{2} + \frac{\lambda'_{\Phi_3\Phi_{B-L}} v_{B-L}^2}{2} + \frac{Y v_{B-L}^2 v_{12}}{8 v_3} \\
M_{\phi_{B-L}}^2 &\simeq 2\lambda_{\Phi_{B-L}} v_{B-L}^2 + \frac{\lambda_{H\Phi_{B-L}} v^2}{2} + \frac{\lambda_{\Phi_{12}\Phi_{B-L}} v_{12}^2}{2} \\
&+ \frac{\lambda'_{\Phi_{12}\Phi_{B-L}} v_{12}^2}{2} + \frac{\lambda_{\Phi_3\Phi_{B-L}} v_3^2}{2} + \frac{\lambda'_{\Phi_3\Phi_{B-L}} v_3^2}{2} + \frac{Y v_{12}^2 v_3}{2}
\end{aligned} \tag{12}$$

where we have neglected the higher order terms involving u_Δ and u_ξ as it is very small and no significant contribution comes from these terms.

The Z_{B-L} boson acquires mass through the vevs of Φ_{B-L}, Φ_{12}, Φ_3 which are charged under $U(1)_{B-L}$ and is given by:

$$M^2_{Z_{B-L}} = g^2_{B-L}(v^2_{B-L} + 64v^2_{12} + 100v^2_3). \tag{13}$$

For simplicity as well as from interesting phenomenological perspective, we assume that the masses of h, ϕ_3 and ξ are of similar order in sub-TeV range, while the masses of ϕ_{B-L} and ϕ_{12} are in a few TeV scale. Thus the assumption for mass heirarchy among scalars is $M_\delta >> M_{\phi_{B-L}}, M_{\phi_{12}} > M_h, M_{\phi_3}, M_\xi$. Now we examine the mixing among the CP-even scalars h, ϕ_3 and ξ. The mixing between h and ξ is highly suppressed as we discussed in section III A and hence negligible. Similarly ξ-Φ_3 mixing is of the order $\mathrm{O}(\frac{u_\xi}{v_3})$ and hence negligibly small. Hence in what follows we consider H and Φ_3 mixing while discussing low energy phenomenology.

Minimising the scalar potential 6 with respect to $\langle H \rangle = v$ and $\langle \Phi_3 \rangle = v_3$, we obtain:

$$v_3 = \sqrt{\frac{2\lambda_{H\Phi_3}M_H^2 - 4\lambda_H M^2_{\Phi_3}}{4\lambda_H\lambda_{\Phi_3} - \lambda^2_{H\Phi_3}}} \quad \text{and} \quad v = \sqrt{\frac{2\lambda_{H\Phi_3}M^2_{\Phi_3} - 4\lambda_{\Phi_3} M^2_H}{4\lambda_H\lambda_{\Phi_3} - \lambda^2_{H\Phi_3}}}. \tag{14}$$

The mass matrix of h and ϕ_3 can be written as:

$$\mathcal{M}^2(h, \phi_3) = \begin{pmatrix} \lambda_H v^2 & \lambda_{H\Phi_3} v v_3 \\ \lambda_{H\Phi_3} v v_3 & \lambda_{\Phi_3} v_3^2 \end{pmatrix}. \tag{15}$$

Here the mixing between h and ϕ_3 is governed by the coupling $\lambda_{H\Phi_3}$. The masses of the physical Higgses can be obtained by diagonalising the above mass matrix as:

$$\begin{aligned} M^2_{h_1} &= \frac{1}{2}[(\lambda_H v^2 + \lambda_{\Phi_3} v_3^2) - \sqrt{(\lambda_{\Phi_3} v_3^2 - \lambda_H v^2)^2 + 4(\lambda_{H\Phi_3} v v_3)^2}] \\ M^2_{h_2} &= \frac{1}{2}[(\lambda_H v^2 + \lambda_{\Phi_3} v_3^2) + \sqrt{(\lambda_{\Phi_3} v_3^2 - \lambda_H v^2)^2 + 4(\lambda_{H\Phi_3} v v_3)^2}] \end{aligned} \tag{16}$$

where h_1 and h_2 can be realised as SM-like Higgs and the second Higgs, respectively. Thus the mass eigen states of these scalars can be given as:

$$\begin{aligned} h_1 &= \cos\beta \; h + \sin\beta \; \phi_3 \\ h_2 &= -\sin\beta \; h + \cos\beta \; \phi_3 \,, \end{aligned} \tag{17}$$

where

$$\tan 2\beta = \Big(\frac{2\lambda_{H\Phi_3} v v_3}{\lambda_{\Phi_3} v_3^2 - \lambda_H v^2}\Big). \tag{18}$$

Thus it is evident that the mixing angle β vanishes if $\lambda_{H\Phi_3} \to 0$ or if $v_3 >> v$.

IV. DARK MATTER

At a TeV scale, the $U(1)_{\rm B-L}$ gauge symmetry breaks down to a remnant Z_2 by the vev of Φ_{12}, $\Phi_{\rm B-L}$ and Φ_3. We assume that the right-handed neutrinos are odd under the Z_2 symmetry, while all other particles even. As a result the lightest right-handed neutrino is a viable candidate of dark matter.

A. Right Handed Neutrinos and Their Interactions

From Eqn. (5), the mass matrix of right handed neutrinos in the effective theory can be given as:

$$-\mathscr{L}^{mass}_{\nu_R} = \frac{1}{2}\begin{pmatrix}\overline{(\nu_{_{R1}})^c} & \overline{(\nu_{_{R2}})^c} & \overline{(\nu_{_{R3}})^c}\end{pmatrix}\mathcal{M}\begin{pmatrix}\nu_{_{R1}}\\ \nu_{_{R2}}\\ \nu_{_{R3}}\end{pmatrix} \tag{19}$$

where

$$\mathcal{M} = \begin{pmatrix} Y_{11}v_{_{12}} & Y_{12}v_{_{12}} & Y_{13}v_{_{B-L}} \\ Y_{12}v_{_{12}} & Y_{22}v_{_{12}} & Y_{23}v_{_{B-L}} \\ Y_{13}v_{_{B-L}} & Y_{23}v_{_{B-L}} & Y_{33}v_3 \end{pmatrix} = \begin{pmatrix} [M_{12}] & [M'] \\ [M']^T & M_3 \end{pmatrix}. \tag{20}$$

Here M_{12}, M', M_3 are:

$$M_{12} = \begin{pmatrix} Y_{11}v_{_{12}} & Y_{12}v_{_{12}} \\ Y_{12}v_{_{12}} & Y_{22}v_{_{12}} \end{pmatrix}, \; M' = \begin{pmatrix} Y_{13}v_{_{B-L}} \\ Y_{23}v_{_{B-L}} \end{pmatrix}, \; M_3 = \begin{pmatrix} Y_{33}v_3 \end{pmatrix}.$$

The above right handed neutrino Majorana mass matrix $\mathcal{M}$ can be block-diagonalised using a rotation matrix of the form :

$$\mathcal{R}_\theta = \begin{pmatrix} \cos\theta \begin{bmatrix} 1 & 0 \\ 0 & 1 \end{bmatrix} & \sin\theta \begin{bmatrix} 1 \\ 1 \end{bmatrix} \\ -\sin\theta \begin{bmatrix} 1 & 1 \end{bmatrix} & \cos\theta \end{pmatrix} = \begin{pmatrix} \cos\theta & 0 & \sin\theta \\ 0 & \cos\theta & \sin\theta \\ -\sin\theta & -\sin\theta & \cos\theta \end{pmatrix} \tag{21}$$

where it can be checked that this transformation is orthogonal in the linear over θ approximation as,

$$(\mathcal{R}_\theta)^T\mathcal{R}_\theta = \begin{pmatrix} 1 & \sin^2\theta & 0 \\ \sin^2\theta & 1 & 0 \\ 0 & 0 & 1+\sin^2\theta \end{pmatrix} \simeq I_3. \tag{22}$$

Here θ is essentially the mixing between M_{12} and M_3. This mixing angle can be expressed as:

$$\sin 2\theta \simeq \frac{Y_{13} v_{B-L}}{Y_{11} v_{12} + Y_{12} v_{12} - Y_{33} v_3} \simeq \frac{Y_{23} v_{B-L}}{Y_{11} v_{12} + Y_{12} v_{12} - Y_{33} v_3}. \tag{23}$$

Then we can completely diagonalise the block-diagonalised matrix $\mathcal{M}$ further by performing another orthogonal transformation using the transformation matrix

$$\mathcal{R}_\alpha = \begin{pmatrix} \cos\alpha & \sin\alpha & 0 \\ -\sin\alpha & \cos\alpha & 0 \\ 0 & 0 & 1 \end{pmatrix}, \tag{24}$$

where the mixing angle α is given by

$$\tan 2\alpha \simeq \frac{Y_{12} v_{12}}{Y_{11} v_{12} - Y_{22} v_{12}}. \tag{25}$$

Thus the mass eigen states of the right handed neutrinos can be written as:

$$\begin{pmatrix} N_{1R} \\ N_{2R} \\ N_{3R} \end{pmatrix} = \begin{pmatrix} \cos\alpha & \sin\alpha & 0 \\ -\sin\alpha & \cos\alpha & 0 \\ 0 & 0 & 1 \end{pmatrix} \begin{pmatrix} \cos\theta & 0 & \sin\theta \\ 0 & \cos\theta & \sin\theta \\ -\sin\theta & -\sin\theta & \cos\theta \end{pmatrix} \begin{pmatrix} \nu_{_{R1}} \\ \nu_{_{R2}} \\ \nu_{_{R3}} \end{pmatrix} \tag{26}$$

Here it is worth noting that the matrix $\mathcal{R} = \mathcal{R}_\alpha \mathcal{R}_\theta$ is also orthogonal ($i.e.$, $\mathcal{R}^T \mathcal{R} = I$) in the linear over θ approximation.

For simplicity we assume that there is strong hierarchy between ν_{R3} and ν_{R1}, ν_{R2}. This implies that $\sin\theta \to 0$ in Eq. 23. In fact, this can be achieved if we assume $Y_{13} v_{\rm B-L}, Y_{23} v_{\rm B-L} << (Y_{11} + Y_{12}) v_{12}$. In this limit, $\nu_{_{R3}}$ completely decouples from $\nu_{_{R1}}$ and $\nu_{_{R2}}$. As a result the diagonalisation of above mass matrix gives the mass eigen values corresponding to the states N_{1R},N_{2R} and N_{3R} as

$$\begin{aligned} M_{1,2} &= \frac{1}{2}[(Y_{11} v_{12} + Y_{22} v_{12}) \pm \sqrt{(Y_{11} v_{12} - Y_{22} v_{12})^2 + 4(Y_{12} v_{12})^2}] \\ M_3 &= Y_{33}\, v_3. \end{aligned} \tag{27}$$

Interactions

The interaction terms of the right handed neutrinos with Z_{B-L} in the mass eigen basis can be written as :

$$\begin{aligned}
\mathcal{L}_{Z_{B-L}} = g_{_{B-L}}\Bigg[&\{(4\cos^2\theta - 5\sin^2\theta)(1+\sin 2\alpha)\}\overline{N_{1R}}\gamma^\mu N_{1R} + \{16\sin^2\theta - 5\cos^2\theta\}\overline{N_{3R}}\gamma^\mu N_{3R} \\
&+ \{(4\cos^2\theta - 5\sin^2\theta)(1-\sin 2\alpha)\}\overline{N_{2R}}\gamma^\mu N_{2R} \\
&+ \cos 2\alpha\{4\cos^2\theta - 5\sin^2\theta\}\Big(\overline{N_{1R}}\gamma^\mu N_{2R} + \overline{N_{2R}}\gamma^\mu N_{1R}\Big) \\
&- \frac{13}{2}\sin 2\theta\{\cos\alpha + \sin\alpha\}\Big(\overline{N_{1R}}\gamma^\mu N_{3R} + \overline{N_{3R}}\gamma^\mu N_{1R}\Big) \\
&- \frac{13}{2}\sin 2\theta\{\cos\alpha - \sin\alpha\}\Big(\overline{N_{2R}}\gamma^\mu N_{3R} + \overline{N_{3R}}\gamma^\mu N_{2R}\Big)\Bigg](Z_{B-L})_\mu. \qquad (28)
\end{aligned}$$

The Yukawa interactions in the mass basis can be written as:

$$\begin{aligned}
Y_{ab}\overline{(\nu_{_{Ra}})^c}\Phi\nu_{_{Rb}} &= Y_{11}\overline{(\nu_{_{R1}})^c}\Phi\nu_{_{R1}} + Y^{22}\overline{(\nu_{_{R2}})^c}\Phi\nu_{_{R2}} + Y^{12}\overline{(\nu_{_{R1}})^c}\Phi\nu_{_{Rb}} + Y^{12}\overline{(\nu_{_{R2}})^c}\Phi\nu_{_{R1}} \\
&= Y\Bigg[\cos^2\theta(1+\sin 2\alpha)\overline{(N_{1R})^c}\phi N_{1R} + \cos^2\theta(1-\sin 2\alpha)\overline{(N_{2R})^c}\phi N_{2R} \\
&+ 4\sin^2\theta\overline{(N_{3R})^c}\phi N_{3R} + \cos^2\theta\cos 2\alpha\Big(\overline{(N_{1R})^c}\phi N_{2R} + \overline{(N_{2R})^c}\phi N_{1R}\Big) \\
&- \sin 2\theta(\cos\alpha + \sin\alpha)\Big(\overline{(N_{1R})^c}\phi N_{3R} + \overline{(N_{3R})^c}\phi N_{1R}\Big) \\
&- \sin 2\theta(\cos\alpha - \sin\alpha)\Big(\overline{(N_{2R})^c}\phi N_{3R} + \overline{(N_{3R})^c}\phi N_{2R}\Big)\Bigg] \qquad (29)
\end{aligned}$$

Similarly

$$\begin{aligned}
Y_{a3}\overline{(\nu_{_{Ra}})^c}\Phi_{B-L}\nu_{_{R3}} &= Y\left[\overline{(\nu_{_{R1}})^c}\Phi_{B-L}\nu_{_{R3}} + \overline{(\nu_{_{R2}})^c}\Phi_{B-L}\nu_{_{R3}} + h.c\right] \\
&= Y\Bigg[\sin 2\theta(1+\sin 2\alpha)\overline{(N_{1R})^c}\phi_{B-L}N_{1R} + \sin 2\theta(1-\sin 2\alpha)\overline{(N_{2R})^c}\phi_{B-L}N_{2R} \\
&- 2\sin 2\theta\overline{(N_{3R})^c}\phi_{B-L}N_{3R} \\
&+ \frac{1}{2}\sin 2\theta(\cos 2\alpha + \sin 2\alpha)\Big(\overline{(N_{1R})^c}\phi_{B-L}N_{2R} + \overline{(N_{2R})^c}\phi_{B-L}N_{1R}\Big) \\
&+ (\cos^2\theta - 2\sin^2\theta)(\cos\alpha - \sin\alpha)\Big(\overline{(N_{2R})^c}\phi_{B-L}N_{3R} + \overline{(N_{3R})^c}\phi_{B-L}N_{2R}\Big) \\
&+ (\cos^2\theta - 2\sin^2\theta)(\cos\alpha + \sin\alpha)\Big(\overline{(N_{2R})^c}\phi_{B-L}N_{3R} + \overline{(N_{3R})^c}\phi_{B-L}N_{2R}\Big)\Bigg]. \qquad (30)
\end{aligned}$$

and

$$Y_{33}\overline{(\nu_{R3})^c}\Phi_3\nu_{R3} = Y_{33}\Big[\sin^2\theta(1+\sin 2\alpha)\overline{(N_{1R})^c}\phi_3 N_{1R} + \sin^2\theta(1-\sin 2\alpha)\overline{(N_{2R})^c}\phi_3 N_{2R}$$
$$+ \cos^2\theta\overline{(N_{3R})^c}\phi_3 N_{3R} + \sin^2\theta\cos 2\alpha\Big(\overline{(N_{1R})^c}\phi_3 N_{2R} + \overline{(N_{2R})^c}\phi_3 N_{1R}\Big)$$
$$+ \frac{1}{2}\sin 2\theta(\cos\alpha+\sin\alpha)\Big(\overline{(N_{1R})^c}\phi_3 N_{3R} + \overline{(N_{3R})^c}\phi_3 N_{1R}\Big)$$
$$+ \frac{1}{2}\sin 2\theta(\cos\alpha-\sin\alpha)\Big(\overline{(N_{2R})^c}\phi_3 N_{3R} + \overline{(N_{3R})^c}\phi_3 N_{2R}\Big)\Big]. \quad (31)$$

But as ϕ_3 mixes with SM Higgs h, then the Eq. 31 can be written in terms of the scalar mass eigen states as:

$$Y_{33}\overline{(\nu_{R3})^c}\Phi_3\nu_{R3} = Y_{33}\sin\beta\Big[\sin^2\theta(1+\sin 2\alpha)\overline{(N_{1R})^c}h_1 N_{1R} + \sin^2\theta(1-\sin 2\alpha)\overline{(N_{2R})^c}h_1 N_{2R}$$
$$+ \cos^2\theta\overline{(N_{3R})^c}h_1 N_{3R} + \sin^2\theta\cos 2\alpha\Big(\overline{(N_{1R})^c}h_1 N_{2R} + \overline{(N_{2R})^c}h_1 N_{1R}\Big)$$
$$+ \frac{1}{2}\sin 2\theta(\cos\alpha+\sin\alpha)\Big(\overline{(N_{1R})^c}h_1 N_{3R} + \overline{(N_{3R})^c}h_1 N_{1R}\Big)$$
$$+ \frac{1}{2}\sin 2\theta(\cos\alpha-\sin\alpha)\Big(\overline{(N_{2R})^c}h_1 N_{3R} + \overline{(N_{3R})^c}h_1 N_{2R}\Big)\Big]$$
$$+ Y^{33}\cos\beta\Big[\sin^2\theta(1+\sin 2\alpha)\overline{(N_{1R})^c}h_2 N_{1R} + \sin^2\theta(1-\sin 2\alpha)\overline{(N_{2R})^c}h_2 N_{2R}$$
$$+ \cos^2\theta\overline{(N_{3R})^c}h_2 N_{3R} + \sin^2\theta\cos 2\alpha\Big(\overline{(N_{1R})^c}h_2 N_{2R} + \overline{(N_{2R})^c}h_2 N_{1R}\Big)$$
$$+ \frac{1}{2}\sin 2\theta(\cos\alpha+\sin\alpha)\Big(\overline{(N_{1R})^c}h_2 N_{3R} + \overline{(N_{3R})^c}h_2 N_{1R}\Big)$$
$$+ \frac{1}{2}\sin 2\theta(\cos\alpha-\sin\alpha)\Big(\overline{(N_{2R})^c}h_2 N_{3R} + \overline{(N_{3R})^c}h_2 N_{2R}\Big)\Big]. \quad (32)$$

B. Relic Abundance of DM

As discussed above, the lightest right-handed neutrino is stable due to a remnant Z_2 symmetry. Without loss of generality we assume $N_3 = (N_{3R} + N_{3R}^c)/\sqrt{2}$ is the lightest right handed neutrino and acts as a candidate of DM.

The relic abundance of DM (N_3) can be achieved from the annihilation channels of N_3 to the SM particles. The various processes which are contributing to the DM relic abundance via the annihilation channels can be seen in Fig. 2.

The relic density of DM is calculated using micrOMEGAs [40]. The masses for the heavier Z_2 odd particles are chosen to be 1.5 TeV and 2 TeV. In Fig. 3, we show the relic density of DM as a function of DM mass. The resonance regions are corresponding to the SM-like Higgs h_1, second

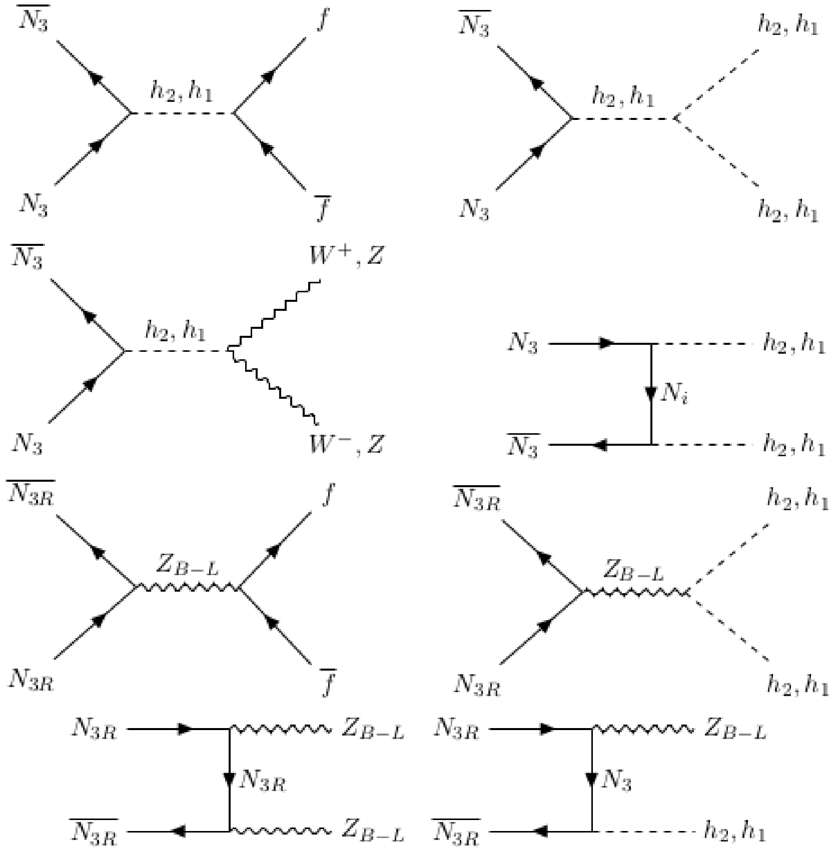

FIG. 2: DM annihilation channels.

Higgs h_2 and gauge boson Z_{B-L}. The masses of second Higgs M_{h_2} and the gauge boson $M_{Z_{B-L}}$ are taken to be 400 GeV and 2000 GeV respectively and the gauge coupling is taken to be $g_{B-L} = 0.035$, where these value are in agreement with the collider constraints [41]. In Fig. 3, an arbitrary choice of set $(Y_{33}^{eff},\ \sin\beta)$ is taken $(0.3, 0.01), (0.01, 0.1), (0.001, 0.5)$. One can find that the relic abundance criteria of DM is typically satisfied at the resonance regions only. The effective coupling Y_{33}^{eff} of the vertex $\overline{(\nu_{R3})^c} h_2 \nu_{R3}$ and $\sin\beta$ are taken to be random numbers between $10^{-3} - 1$ and the points which are allowed for Y_{33}^{eff} and $\sin\beta$ by the relic density constraint are shown in Fig. 8 (Orange points). We also show the allowed parameter space in the $g_{B-L} - M_{Z_{B-L}}$ plane which are allowed by the relic density constraints for three different values of DM mass in Fig 4. Fig 5 shows the plots of g_{B-L} vs M_{N_3} which satisfies the observed relic density for $M_{Z_{B-L}} = 2$ TeV, 3 TeV and 4 TeV respectively.

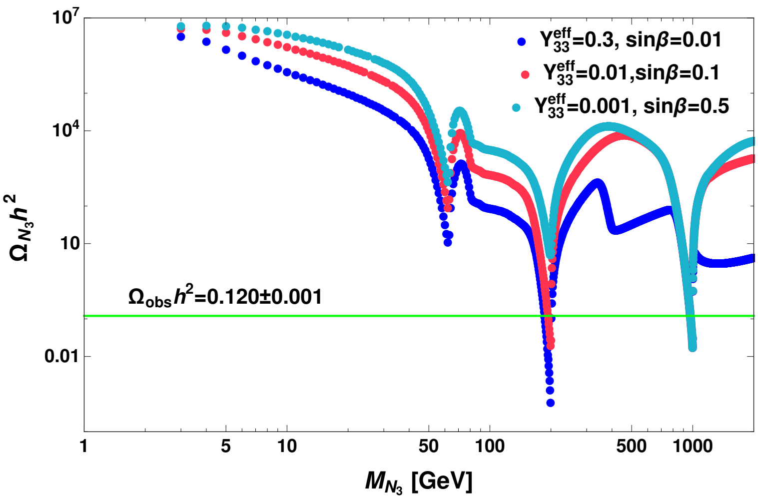

FIG. 3: Relic density of DM as a function of its mass. Contribution to the relic density of DM is only through its annihilation channels. The Green horizontal line shows the observed relic density of DM [4].

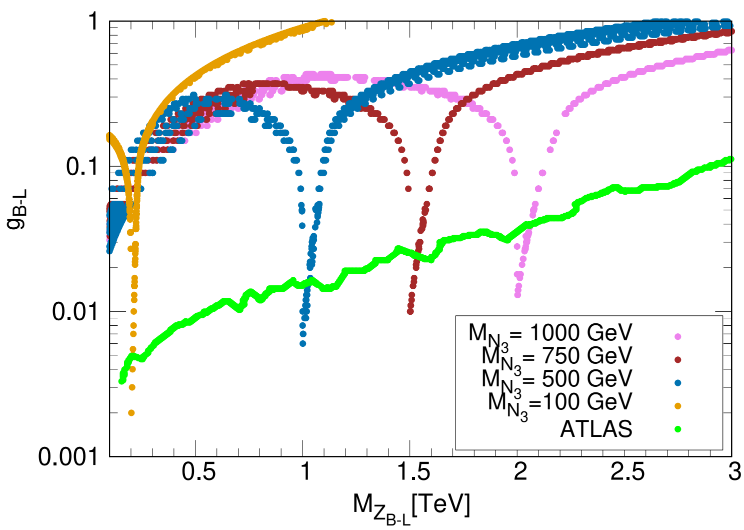

FIG. 4: Points allowed by the relic density constraint in the plane of $M_{Z_{B-L}}$ vs g_{B-L}. The dips corresponds to the resonance for the mass of Z_{B-L}, *i.e.*, $M_{Z_{B-L}}$=200(Orange), 1000(Blue), 1500(Maroon), 2000(Pink) GeV, corresponding to the DM mass 100, 500, 750 and 1000 GeV respectively. The Geen line shows the ATLAS constraint on effectibe coupling strength g_{B-L} as a function of gauge boson mass [41]. Here the effective coupling Y_{33}^{eff} and $\sin\beta$ are taken to be random numbers between $10^{-3}-1$.

C. Direct Detection

The spin-independent scattering of DM is possible through $\phi_3 - h$ mixing, where DM particles can scatter off the target nuclei which are located at terrestrial laboratories. The corresponding Feynmann diagram can be seen in Fig. 6. The spin-independent elastic scattering cross-section of

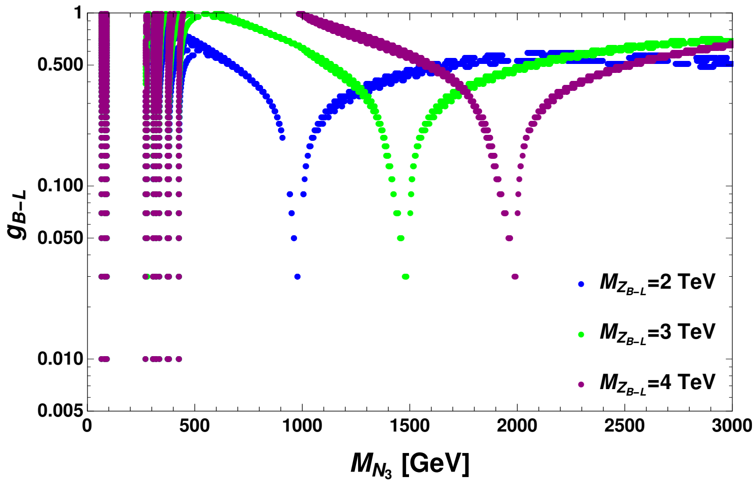

FIG. 5: Points allowed by the relic density constraint in the plane of M_{N_3} vs g_{B-L}. The dips correspond to the resonance corresponding to the mass of Z_{B-L}, *i.e.*, $M_{Z_{B-L}}$=2000(Blue), 3000(Green), 4000(Purple) GeV. The effective coupling Y_{33}^{eff} and $\sin\beta$ are taken to be random numbers between $10^{-3} - 1$.

DM per nucleon can be expressed as

$$\sigma_{SI}^{h_1 h_2} = \frac{\mu_r^2}{\pi A^2}\left[Z f_p + (A-Z) f_n\right]^2 \tag{33}$$

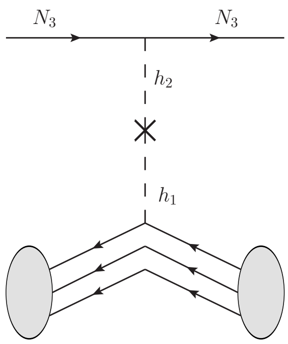

FIG. 6: The spin-independent scattering cross-section of DM-nucleon via Higgs portal.

where $\mu_r = M_{N_3} m_n/(M_{N_3} + m_n)$ is the reduced mass, where m_n is the nucleon (proton or neutron) mass. The A and Z are the mass and atomic number of the target nucleus, respectively. The f_p and f_n are the interaction strengths of proton and neutron with DM, respectively and they can be given as,

$$f_{p,n} = \sum_{q=u,d,s} f_{T_q}^{p,n} \alpha_q \frac{m_{p,n}}{m_q} + \frac{2}{27} f_{TG}^{p,n} \sum_{q=c,t,b} \alpha_q \frac{m_{p,n}}{m_q}\,, \tag{34}$$

where

$$\alpha_q = \frac{Y_{33}^{eff} \sin 2\beta}{2\sqrt{2}} \left(\frac{m_q}{v}\right)\left[\frac{1}{M_{h_2}^2} - \frac{1}{M_{h_1}^2}\right]. \tag{35}$$

In the above Eq. 34, the values of $f^{p,n}_{T_q}$ can be found in [42].

Using Eq. 34 and 35, the spin-independent cross-section in Eq. 33, can be re-expressed as:

$$\sigma_{\rm SI}^{h_1 h_2} = \frac{{\mu_r}^2}{\pi A^2}\left(\frac{Y_{33}^{eff}\sin 2\beta}{2\sqrt{2}}\right)^2\left[\frac{1}{M_{h_2}^2}-\frac{1}{M_{h_1}^2}\right]^2$$
$$\times\left[Z\left(\frac{m_p}{v}\right)\left(f_{Tu}^p+f_{Td}^p+f_{Ts}^p+\frac{2}{9}f_{TG}^p\right)+(A-Z)\left(\frac{m_n}{v}\right)\left(f_{Tu}^n+f_{Td}^n+f_{Ts}^n+\frac{2}{9}f_{TG}^n\right)\right]^2. \quad (36)$$

The more stringent bound on direct detection come from the XENON1T [43], where it rules out spin-independent cross-section down to $\sigma_{\rm SI} \approx 10^{-47}{\rm cm}^2$. The Fig. 7, show the parameter space which is allowed by the XENON1T bound. The only free parameters in Eq. 36 are λ_{33}^{eff} and $\sin\beta$ These values are varied randomly between $10^{-3}-1$. The parameter space allowed by λ_{33}^{eff} and $\sin\beta$ are shown in Fig. 8, with Sky blue points, which keeps the spin-independent DM-nucleon cross-section below the XENON1T bound. The Orange points are allowed by the relic density bound. The overlapped region are the points which are allowed by the both relic density and XENON1T bounds.

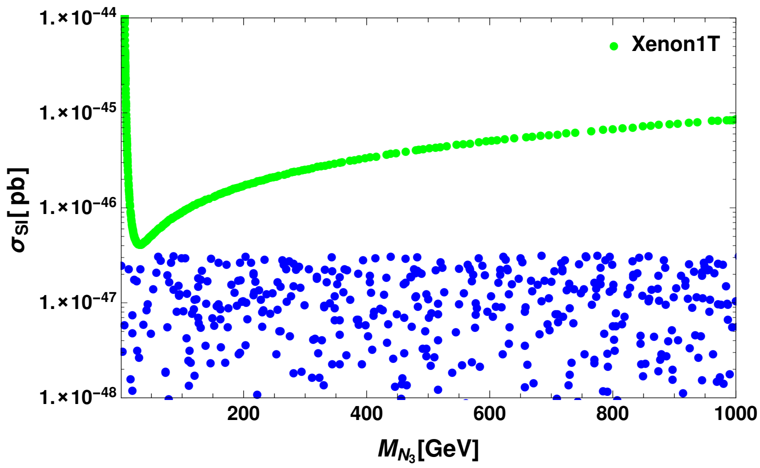

FIG. 7: The spin-independent DM-nucleon cross-section (Blue dots) which is not yet ruled out by XENON1T bound (Green dots).

V. COLLIDER SIGNATURES

As the model is extended with $U(1)_{B-L}$ gauge group and all the fermions as well as most of the scalars like ξ,Φ_{B-L},Φ_{12},Φ_3 are charged under $B-L$ gauge interaction, so we can encounter interesting collider signatures predicted by this model at LHC or FCC [21, 44, 45]

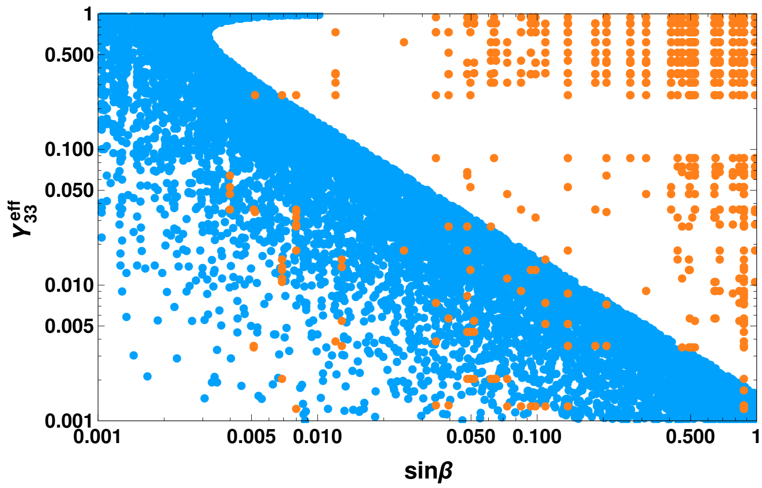

FIG. 8: The parameter region allowed by relic density and XENON1T bounds. The Orange points are allowed by the relic density constraint. The Sky blue points allowed by the XENON1T bound.

If $M_{Z_{B-L}} < M_\xi$,$M_{\Phi_{B-L}}$,$M_{\Phi_{12}}$,M_{Φ_3}, then as the gauge coupling is proportional to the $B-L$ charges, so Z_{B-L} can dominantly decay to the lightest and next to lightest stable particles of the dark sector i.e $Z_{B-L} \to N_{3R}\overline{N_{3R}}$ or $Z_{B-L} \to N_{iR}\overline{N_{iR}}$ with i=1,2 , (since they have $B-L$ charges $+5$ and -4 respectively). Also decay of Z_{B-L} to a pair of SM leptons (having $B-L$ quantum number -1) will be dominant as compared to it's decay to quarks ($B-L$ quantum number 1/3). But alternatively if $M_{Z_{B-L}} > M_\xi$,$M_{\Phi_{B-L}}$,$M_{\Phi_{12}}$,M_{Φ_3} then the total decay width of Z_{B-L} significantly increases as it additionally decays to $\phi_3\phi_3^*$, $\phi_{12}\phi_{12}^*$, $\xi^{\pm\pm}\xi^{\mp\mp}$, $\xi^\pm\xi^\mp$, $\xi^0\xi^{0*}$ and $\phi_{B-L}\phi_{B-L}^*$.

As already mentioned, this model not only contains the TeV scale $B-L$ gauge boson but also has a triplet scalar ξ with $M_\xi \lesssim 1$TeV. So depending on the relative magnitudes of $M_{Z_{B-L}}$ and $M_{\xi^{\pm\pm}}$, the production cross-section of $\xi^{\pm\pm}$ and Z_{B-L} will vary. The collider search for this triplet scalar in various models have been studied in many papers [25–30].

If $M_{Z_{B-L}} > 2M_{\xi^{\pm\pm}}$, then at LHC $\xi^{\pm\pm}$ particles can be pair produced via Z_{B-L} decay with a significant branching fraction. But if $M_{Z_{B-L}} < 2M_{\xi^{\pm\pm}}$, then at LHC $\xi^{\pm\pm}$ particles can be produced via Drell-Yan process ($q\bar{q} \to \xi^{\pm\pm}\xi^{\mp\mp}$). Once $\xi^{\pm\pm}$ are produced ,then it can decay to different SM particles which can be studied at present and future colliders. The $\xi^{\pm\pm}$ particle can decay to two like-sign charged leptons ($l_\alpha^+ l_\beta^+, \alpha, \beta = e, \mu, \tau$) or to W^+W^+. Though it can also decay to $W^+\xi^+$, but this decay rate [$\Gamma(\xi^{++} \to W^+\xi^+)$] is phase space suppressed since mass of ξ^+ is of the same order of mass of ξ^{++}. It is worth mentioning that for small u_ξ, $\Gamma(\xi^{++} \to l_\alpha^+ l_\beta^+)$ dominates over $\Gamma(\xi^{++} \to W^+W^+)$ as $\Gamma(\xi^{++} \to l_\alpha^+ l_\beta^+)$ varies inversely with u_ξ^2 but $\Gamma(\xi^{++} \to W^+W^+)$ varies directly

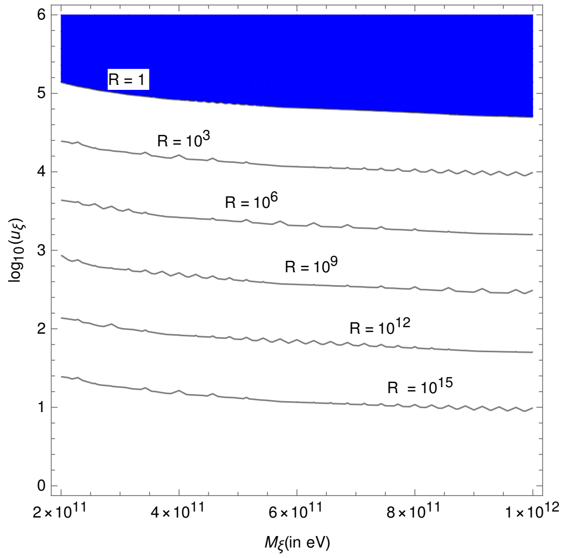

FIG. 9: The contour for R in the plane of $log_{10}(u_\xi)$ and M_ξ.

with u_ξ^2. This is evident from the partial decay widths which are given as:

$$\Gamma(\xi^{++} \to l_\alpha^+ l_\beta^+) = \frac{M_{\xi^{++}}}{4\pi u_\xi^2 (1+\delta_{\alpha\beta})} |(M_\nu)_{\alpha\beta}|^2 \tag{37}$$

and

$$\Gamma(\xi^{++} \to W^+W^+) = g^4 u_\xi^2 M_{\xi^{++}}^3 \left[1 - 4\Big(\frac{M_W}{M_{\xi^{++}}}\Big)^2\right]^{\frac{1}{2}} \left[1 - 4\Big(\frac{M_W}{M_{\xi^{++}}}\Big)^2 + \Big(\frac{M_W}{M_{\xi^{++}}}\Big)^4\right] \tag{38}$$

This can be well analysed by plotting contours of the ratio

$$R = \frac{\Gamma(\xi^{++} \to l_\alpha^+ l_\beta^+)}{\Gamma(\xi^{++} \to W^+W^+)} \tag{39}$$

in the plane of $M_{\xi^{++}}$ vs u_ξ as shown in the Fig 9. Only in the region above $R = 1$ contour where $R < 1$ ξ dominantly decays to W^+W^+ while for all other regions the dilepton decay is the dominant one. This like-sign dilepton channel of ξ^{++} is almost background free and can be seen at LHC. The mass of ξ^{++} is approximately about the invariant mass of the two like sign leptons. The small SM background for this dilepton signature coming from two Z-boson decay can be removed by proper selection of cuts.

The singly charged scalar particles $\xi^\pm$ can be produced along with the doubly charged scalars through the Drell-Yan process ($q\overline{q'} \to \xi^{\mp\mp}\xi^\pm$) mediated via the charged weak gauge boson $W^{\pm *}$. Also there can be pair production of $\xi^\pm$ similar to $\xi^{\pm\pm}$ through Drell-Yan process ($q\bar{q} \to \xi^\mp\xi^\pm$) mediated by γ^*,Z^*,Z_{B-L}^*. Once produced, the decay of $\xi^\pm$ can dominantly ocuur through the channel $\xi^+ \to l^+ + \nu$. But since the neutrinos are invisible at the detector, the decay of $\xi^\pm$, produced through the channel $q\overline{q'} \to \xi^{\mp\mp}\xi^{\pm}$ mediated via the the charged weak gauge boson $W^{\pm *}$, will lead to a three lepton final state ($l^{\pm}l^{\pm}l^{\mp}$). On the other hand, the decay of $\xi^{\pm}$, produced through the channel $q\bar{q} \to \xi^{\mp}\xi^{\pm}$ will lead to a two lepton final state ($l^{\pm}l^{\mp}$). However in both the cases we will have large SM background, so with proper application of cuts, these events can be studied at LHC.

VI. CONCLUSION

In this paper we studied a gauged $U(1)_{B-L}$ extension of a TeV scale type-II seesaw. We implemented it by introducing two scalar triplets Δ ($M_{\Delta} \simeq 10^{14}$ GeV) and ξ ($M_{\xi} \leq 1$ TeV) with $M_{\xi} << M_{\Delta}$. Even though there is orders of magnitude difference between the masses of these two scalars but their contribution to the neutrino mass is identical because of the small mixing between them which arises at TeV scale when $U(1)_{B-L}$ is broken by the vev of the singlet scalar Φ_{B-L}. Δ being super heavy is decoupled from the low energy effective theory however the decay of leptophillic ξ^{++} is almost background free and can be studied at LHC. To make the theory free from anomalies we introduced three right handed neutrino fields $\nu_{R_i}(i = 1, 2, 3)$ which have charges under $U(1)_{B-L}$ as -4, -4, +5 respectively. The $U(1)_{B-L}$ charges of right handed neutrinos precludes any coupling with the SM fermions. The lightest one among these right handed neutrinos is the DM candidate of the model whose stability is guaranteed by a remnant Z_2 symmetry after $U(1)_{B-L}$ breaking. We showed a combined parameter space which allows both the relic and direct detection bounds. The model under consideration here explains the smallness of neutrino masses as well as DM. To satisfy the relic density constraint of the DM we have considered only the contribution coming from annihilation cross-sections of the DM. We will study the contribution of co-annihilation channels in a future project.

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